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On the Nature of Precursors in the Radio Pulsar Profiles

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ABSTRACT

In the average profiles of several radio pulsars, the main pulse is accompanied by the preceding component. This so called precursor is known for its distinctive polarization, spectral, and fluctuation properties. Recent single-pulse observations hint that the sporadic activity at the extreme leading edge of the pulse may be prevalent in pulsars. We for the first time propose a physical mechanism of this phenomenon. It is based on the induced scattering of the main pulse radiation into the background. We show that the scattered component is directed approximately along the ambient magnetic field and, because of rotational aberration in the scattering region, appears in the pulse profile as a precursor to the main pulse. Our model naturally explains high linear polarization of the precursor emission, its spectral and fluctuation peculiarities as well as suggests a specific connection between the precursor and the main pulse at widely spaced frequencies. This is believed to stimulate multifrequency single-pulse studies of intensity modulation in different pulsars.

Key words: pulsars: general – radiation mechanisms: non-thermal – scattering

1 INTRODUCTION

The radio profile of a pulsar may include the precursor - a peculiar component preceding the main pulse by up to a few tens degrees in longitude. The presence of such a component is more or less firmly ascertained for the Crab pulsar (Campbell, Heiles & Rankin 1970), PSR B1055-52 (McCulloch et al. 1976), PSR B1822-09 (Fowler, Wright & Morris 1981), and the Vela pulsar (Krishnamohan & Downs 1983). The precursor differs substantially from the components of the main pulse by its spectral and polarization properties. In the main pulse, the component widths and separations generally increase with wavelength, signifying the overall broadening of the emission cone. At the same time, the width of the precursor and its separation from the main pulse remains unchanged over a broad frequency range. The spectrum of the precursor is also distinct. In the pulsar B1822-09, it is unusually flat (Fowler et al. 1981; Gil et al. 1994), whereas in the Crab pulsar it is extremely steep (Manchester 1971). The most distinctive feature of the precursor is its complete linear polarization, in contrast to the main pulse emission, which is typically somewhat depolarized due to simultaneous presence of the two orthogonally polarized modes.

The single-pulse studies have further revealed some fascinating features of the precursor emission. The precursor intensity shows strong pulse-to-pulse fluctuations. In the Vela pulsar, the intensity variations affect the longitudinal location of the precursor: the stronger the precursor, the larger is its separation from the main pulse (Krishnamohan & Downs 1983). In PSR B1822-09, the precursor appears pronounced only in strong enough pulses; weak pulses, instead, include the interpulse component, which lags the main pulse by about 180° in longitude (Fowler et al. 1981; Fowler & Wright 1982; Gil et al. 1994).

Because of sporadic character of the precursor emission, it does not necessarily form a component in the average profile. For example, the pulsars B0950+08 and B1656+14 do not exhibit a precursor in the average profile, but occasional pulses do incorporate strong subpulses at the extreme leading edge of the pulse (Hankins & Cordes 1981; Weltevrede et al. 2006). Thus, in a broad sense, the precursor phenomenon can be defined as a sporadic activity in the longitudinal region preceding the main pulse. One can expect that the precursors defined in this way are much more typical of pulsars than is usually believed.

It is interesting to note that the precursor phenomenon appears related to other manifestations of intensity modulation in pulsar radio emission. Recent single-pulse observations of PSR J1326-6700 (Wang, Manchester & Johnston 2007) reveal the precursor component arising exceptionally during the nulls of the main pulse, when the intensity of the latter drops below the detection level. Furthermore, several other pulsars of the above survey exhibit occasional alternation between the two forms of the average profile (mode changing phenomenon), the modal profiles being mainly different in the intensity of the leading component relatively to that of the rest of the profile.

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Unusual spectral, polarization, and fluctuation properties of the precursors make them puzzling. To the best of our knowledge, no attempts have previously been aimed at explaining the physics of this phenomenon. Recently Dyks, Zhang & Gil (2005) have suggested a geometrical model for the profile of PSR B1822-09. These authors assume that the main pulse and the precursor originate independently at different locations in the magnetosphere and the precursor emission intermittently reverses its direction to form the interpulse. In that model, the mechanism of reversal of the emission direction remains obscure, but for any conceivable switching mechanism it is principally difficult to explain its dependence on the main pulse intensity.

In the present paper, we for the first time propose a physical mechanism of the precursor formation. In our model, the precursor arises as a result of induced scattering of the main pulse emission into the background by the particles of the ultrarelativistic highly magnetized plasma of a pulsar. Our mechanism naturally explains the observed polarization, spectral, and fluctuation properties of the precursor emission as well as suggests its connection to the main pulse.

2 MAGNETIZED INDUCED SCATTERING

2.1 Statement of the problem

Pulsar radio emission is generated deep in the magnetosphere inside of the open field line tube. Hence, it originates and propagates in the flow of the ultrarelativistic electron-positron plasma, which streams along the open magnetic lines. As the brightness temperatures of the radio emission are extremely high, $T_B \sim 10^{25}-10^{30}$ K, the waves may be subject to efficient induced scattering off the plasma particles. According to the radio emission theories based on the plasma instabilities, the frequency of the generated waves is close to the local Lorentz-shifted proper plasma frequency, $\omega\sim\omega_p\sqrt{\gamma}$, where $\omega_p\equiv\sqrt{4\pi n_e e^2/m}$, n_e is the plasma number density, e and m are the electron charge and mass, and γ is the plasma Lorentz-factor (but see Melrose & Gedalin 1999, for the criticism of this point). In the vicinity of the emission region, where the above condition is still valid, induced scattering off the plasma particles is a collective process. The transverse waves are involved in the induced three-wave interactions (Luo & Melrose 2006). A particular case of induced Raman scattering in application to the pulsar magnetosphere has been considered by Gangadhara & Krishan (1993) and Lyutikov (1998).

As the plasma number density decreases with distance from the neutron star along with the magnetic field strength, $n_e\propto B\propto r^{-3}$, well above the emission region $\omega\gg\omega_p\sqrt{\gamma}$ and the collective effects become negligible. The external magnetic field significantly affects the scattering process on condition that the radio wave frequency in the particle rest frame is much less than the electron gyrofrequency, $\omega'\ll\omega_G\equiv eB/mc$. This condition is valid up to the radius of cyclotron resonance, which lies in the outer magnetosphere. In the present paper, we examine the induced scattering, which takes place well above the emission region and well below the cyclotron resonance radius. Then the magnetized induced Compton scattering is a single-particle process and the incident waves are approximately transverse electromagnetic waves polarized either in the plane of the wavevector and the ambient magnetic field or perpendicularly to this plane.

In application to pulsar magnetosphere, induced scattering in a superstrong magnetic field has first been considered by Blandford & Scharlemann (1976) and found to be efficient. Later on the process has been suggested to explain a number of phenomena in pulsar radio emission (Lyubarskii & Petrova 1996; Petrova 2004a,b). In the present paper, we consider the pulsar radio beam scattering into the background and particularly concentrate on the growth of the scattered component, which is identified with the precursor component of the pulse profile.

2.2 Basic equations

In the preceding literature on the induced scattering in a superstrong magnetic field, the kinetic equation for photon occupation numbers is derived from an analysis of the scattering by a single relativistic electron and does not include the particle distribution function explicitly, so that it directly corresponds to the cold plasma case. In the present paper, we generalize the kinetic equation for the more realistic case of a hot plasma. The corresponding formalism has also been outlined in Blandford & Scharlemann (1976), but we shall go through the derivation once more in order to correct a slight error in that paper.

In the approximation of an infinitely strong magnetic field, the scattering cross-section in the electron frame takes the form

$$\frac{\mathrm{d}\sigma}{\mathrm{d}\Omega_1'}=r_e^2\sin^2\theta'\sin^2\theta_1'\,, \tag{1}$$

where r_e is the classical electron radius, θ' and θ_1' are the propagation angles of the incident and scattered photons, respectively, $\mathrm{d}\Omega_1'$ is the elementary solid angle for the scattered photons, and the primes denote the quantities of the electron rest frame. The cross-section (1) corresponds to the scattering between the photons with the ordinary polarization, i.e with the electric vectors lying in the plane of the ambient magnetic field. The scattering involving the extraordinary polarization states is negligible, since any perturbed motion of a particle (in the field of the incident wave) perpendicular to the ambient magnetic field is suppressed. Let us consider the scattering in the laboratory frame between the two photon states, $\boldsymbol{k}$ and $\boldsymbol{k_1}$, involving the electrons with the momenta p and $p+\delta p$ along the magnetic field. In the scattering act, the momentum parallel to the magnetic field is conserved:

$$\delta p=\hbar k\cos\theta-\hbar k_1\cos\theta_1\,. \tag{2}$$

The probability of generating spontaneously scattered photons per electron per unit time is given by

$$\frac{\mathrm{d}P}{\mathrm{d}t}=n(\boldsymbol{k})\frac{\mathrm{d}^3\boldsymbol{k}}{(2\pi)^3}\eta\frac{\mathrm{d}\sigma}{\mathrm{d}\Omega_1}c^4\frac{\mathrm{d}^3\boldsymbol{k_1}}{\omega_1^2}\delta\left(\omega_1-\frac{\eta}{\eta_1}\omega\right)\,, \tag{3}$$

where $n(\boldsymbol{k})$ is the photon occupation number, $\eta\equiv 1-\beta\cos\theta$, β is the particle velocity in units of c, and the argument of the delta-function signifies the equality of the initial and final frequencies in the particle rest frame: $\omega\gamma\eta=\omega_1\gamma\eta_1$.

The cross-section in the laboratory frame is related to equation (1) by means of relativistic transformations, $\mathrm{d}\Omega_1' = \mathrm{d}\Omega_1/(\gamma\eta_1)^2$, $\sin\theta' = \sin\theta/(\gamma\eta)$. The rate of change of the photon occupation number as a result of induced scatterings is

$$\frac{\mathrm{d}n}{\mathrm{d}t}\frac{\mathrm{d}^3\boldsymbol{k}}{(2\pi)^3} = \int [f(p+\delta p) - f(p)]\,\frac{\mathrm{d}P}{\mathrm{d}t} n_1 \mathrm{d}p\,, \tag{4}$$

where $f(p)$ is the electron distribution function. Note that Blandford & Scharlemann (1976) write down literally the same formula, which incorporates the distribution function in Lorentz-factor, $N(\gamma)$ (cf. eq.[21] in that paper), but it is not true. As $p \equiv \beta\gamma mc$, then $N(\gamma) = f(p)\mathrm{d}p/\mathrm{d}\gamma = mcf(p)/\beta$, where β is also an implicit function of p, i.e. $\beta(p+\delta p) \neq \beta(p)$. In the above kinetic equation, $\mathrm{d}/\mathrm{d}t \equiv \partial/\partial t + c\partial/\partial r$. Hereafter we consider the stationary case, omitting the explicit time derivative. Taking into account that $f(p+\delta p) - f(p) \approx \delta p(\partial f/\partial p)$ and integrating equation (4) by parts yield

$$c\frac{\partial n}{\partial r}\frac{\mathrm{d}^3\boldsymbol{k}}{(2\pi)^3} = -\int f(p)\frac{\partial}{\partial p}\left[\delta p\frac{\mathrm{d}P}{\mathrm{d}t}n_1\right]\mathrm{d}p\,. \tag{5}$$

This can be easily rewritten in terms of $N(\gamma)$, and incorporating equations (1)-(3) we obtain

$$\frac{\partial n}{\partial r} = \int \mathrm{d}\gamma N(\gamma)\int \frac{\mathrm{d}\Omega_1 r_e^2 c^4\hbar\beta(\cos\theta_1 - \cos\theta)}{mc^3} \times\frac{\partial}{\partial\gamma}\left[\frac{nn_1\sin^2\theta\sin^2\theta_1\delta(\omega_1\eta_1 - \omega\eta)k_1^2\mathrm{d}k_1}{\gamma^6\eta^2\eta_1^3\omega_1}\right]. \tag{6}$$

The difference of the above equation from eq.[23] in Blandford & Scharlemann (1976) lies in that the factor β is not subject to differentiation with respect to γ. Integration of equation (6) with the help of delta-function and subsequent differentiation yield finally

$$\frac{\partial n}{\partial r} = \frac{\hbar r_e^2 n\sin^2\theta}{mc}\int N\mathrm{d}\gamma\int\left\{\frac{(\eta_1-\eta)^2}{\eta_1^2\gamma^3}\frac{\partial k_1^2 n_1}{\partial k_1} + \frac{6k\eta^2(\eta_1-\eta)n_1}{\gamma\eta_1^2}\left[1 - \frac{\eta_1+\eta}{2\gamma^2\eta\eta_1}\right]\right\}\frac{\sin^2\theta_1\mathrm{d}\Omega_1}{\gamma^6\eta^3\eta_1^3\beta^2}. \tag{7}$$

For the monoenergetic distribution of the particles, this exactly coincides with the kinetic equation obtained by Blandford & Scharlemann (1976) in another way (cf. their eq. [30]). Our equation (7) presents a correct generalization for an arbitrary distribution function of particles. One can see that in general the detailed form of the particle distribution does not play a crucial role, and hence the characteristic features of the scattering studied below are determined by the role of the external magnetic field. Thus, the evolution of the photon occupation numbers may well be described by incorporating the monoenergetic distribution with some characteristic Lorentz-factor. Seeing that the actual distribution function of the plasma in pulsar magnetosphere is still obscure, this is really a proper assumption.

3 RADIO BEAM SCATTERING INTO BACKGROUND

We are interested in the problem of induced scattering of pulsar radio beam into the background. Far from the emission region, the pulsar radiation propagates quasi-transversely with respect to the ambient magnetic field, $1/\gamma \ll \theta_\mathrm{b} < 1$, and it is known to be highly directional, i.e. $\Delta\theta_\mathrm{b} \lesssim 1/\gamma$. Therefore at any point of the scattering region the incident radiation can be represented by a single wavevector $\boldsymbol{k}_\mathbf{b}$. In case of efficient scattering, the orientations of the scattered photons are believed to concentrate near $\theta_\mathrm{bg}^\mathrm{max}$, which corresponds to the maximum scattering probability at a fixed θ_b. (Note that the kinetic equation (7) is independent of the azimuthal coordinates of the photons.)

To specify $\theta_\mathrm{bg}^\mathrm{max}$ we investigate the evolution of the background occupation numbers using equation (7) and prescribing $n(k,\theta)$ to the background and $n_1(k_1,\theta_1)$ to the beam. Given that $\theta_\mathrm{b} \gg 1/\gamma$ the second term on the right-hand side is $\theta_\mathrm{b}^2\gamma^2$ times larger than the first one and decreases with θ_b as $|\theta_\mathrm{b}^2 - \theta_\mathrm{bg}^2|/\theta_\mathrm{b}^4$. Therefore one can expect that the whole integrand peaks at $\theta_\mathrm{bg}^\mathrm{max} \sim 1/\gamma$. Exact numerical calculations of the scattering probability in the range $\theta_\mathrm{bg} \lesssim 1/\gamma$ reveal the following features. The rate of increase of the background occupation numbers does peak at $\theta_\mathrm{bg}^\mathrm{max} \sim 1/\gamma$, the peak width at half maximum is $\lesssim 1/\gamma$, and the first term of equation (7) can still be ignored. In the kinetic equation describing the rate of change of the beam occupation numbers (which is given by equation (7) with $\theta = \theta_\mathrm{b}$ and $\theta_1 = \theta_\mathrm{bg}^\mathrm{max}$), the first term is also negligible as compared to the second one.

As $1/\gamma$ is a small parameter, we examine induced scattering between the two photon states with the fixed polar angles $\theta \gg 1/\gamma$ and $\theta_1 = \theta_\mathrm{bg}^\mathrm{max} \sim 1/\gamma$ and the frequencies related as $\omega_1 = \omega\eta/\eta_1 \sim \omega\theta^2\gamma^2 \gg \omega$. It is convenient to introduce the photon intensities, $i_{\nu_a} = \hbar\omega^3 n(\boldsymbol{k})/2\pi^2c^2$ and $i_{\nu_b} = \hbar\omega_1^3 n_1(\boldsymbol{k}_1)/2\pi^2c^2$, where $\nu_a \equiv \omega/2\pi$ and $\nu_b \equiv \omega_1/2\pi$. The angular distributions have the form of delta-function and one can integrate them over the solid angle to obtain the spectral intensities $I_{\nu_{a,b}} = \int i_{\nu_{a,b}}\mathrm{d}\Omega_{a,b}$. Then the evolution of the beam intensities in the course of induced scattering is approximately described as (see also Petrova 2004b):

$$\frac{\mathrm{d}I_{\nu_a}}{\mathrm{d}r} = -aI_{\nu_a}I_{\nu_b}\,,$$
$$\frac{\mathrm{d}I_{\nu_b}}{\mathrm{d}r} = aI_{\nu_a}I_{\nu_b}\,, \tag{8}$$

where

$$a \sim \frac{24n_e r_e^2}{m\nu_a^2\gamma^5\theta^4} \tag{9}$$

and n_e is the particle number density, $n_e \equiv \int N(\gamma)\mathrm{d}\gamma$. The above system of equations has the first integral,

$$I_{\nu_a} + I_{\nu_b} = \mathrm{const} \equiv I\,. \tag{10}$$

Thus, our approximate consideration leads to the conservation of the total intensity of the two beams. This actually implies that the intensity redistribution between the beams is much more rapid than the change of the total intensity I, and below we concentrate on the photon transfer between the beams. The photons are mainly transferred to the higher-frequency state, $\nu_b \gg \nu_a$, and become almost aligned with the ambient magnetic field, $\theta_1 \sim 1/\gamma$. This contrasts with the non-magnetic scattering, in which case the kinetic equation contains only the term qualitatively similar to the first term in equation (7), the maximum scattering rate is

in the direction antiparallel to the particle velocity, and the photons shift monotonically toward lower frequencies.

The solution of the system (8) is given by

$$I_{\nu_a} = \frac{I(I^{(0)}_{\nu_a}/I^{(0)}_{\nu_b})\exp(-Iar)}{1+(I^{(0)}_{\nu_a}/I^{(0)}_{\nu_b})\exp(-Iar)}\,,$$
$$I_{\nu_b} = \frac{I}{1+(I^{(0)}_{\nu_a}/I^{(0)}_{\nu_b})\exp(-Iar)}\,. \quad (11)$$

One can see that the efficiency of intensity transfer between the beams is determined by $\Gamma \equiv Iar$. At large enough Γ, $I_{\nu_a} \to 0$ and $I_{\nu_b} \to I$. Since the initial intensity of pulsar beam greatly exceeds that of the background radiation, $I^{(0)}_{\nu_b}/I^{(0)}_{\nu_a} \ll 1$, the intensity transfer to the background is significant on condition that

$$\left(I^{(0)}_{\nu_b}/I^{(0)}_{\nu_a}\right)\exp(\Gamma) \gtrsim 1\,. \quad (12)$$

Note that Γ includes the total intensity I, i.e. actually $I^{(0)}_{\nu_a}$, and for large enough Γ the scattering is efficient independently of the smallness of the background intensity.

4 NUMERICAL ESTIMATES

Now we are to estimate the scattering efficiency, $\Gamma = Iar$, where $I \approx I^{(0)}_{\nu_a}$ and a is given by equation (9), for the parameters relevant to pulsar magnetosphere. The spectrum of pulsar radiation generally has the power-law form, $I^{(0)}_{\nu_a} = I_{\nu_0}(\nu_a/\nu_0)^{-\alpha}$, and the spectral intensity at frequencies $\nu_0 \sim 100$ MHz is related to the total radio luminosity of a pulsar, L, as $I_{\nu_0} = L/\nu_0 S$, where $S = \pi r^2 w^2/4$ is the cross-section of the pulsar beam at a distance r, w the pulse width in the angular measure. It is convenient to normalize the number density of the scattering particles to the Goldreich-Julian number density, $n_e = \kappa B/Pce$, where κ is the plasma multiplicity factor, and P is the pulsar period. Keeping in mind that $n_e \propto B \propto r^{-3}$, one can estimate the scattering efficiency as

$$\Gamma = 10P^{-1}\frac{L}{10^{28}\,\mathrm{erg\,s^{-1}}}\frac{B_\star}{10^{12}\,\mathrm{G}}\left(\frac{\nu_a}{10^8\,\mathrm{Hz}}\right)^{-2-\alpha}$$
$$\times\frac{\kappa}{10^2}\frac{10^2}{\gamma}\left(\frac{w}{0.4}\right)^{-2}\left(\frac{\theta}{0.1}\frac{\gamma}{10^2}\frac{r}{10^8\,\mathrm{cm}}\right)^{-4}, \quad (13)$$

where $B_\star$ is the magnetic field strength at the stellar surface, and it is taken that the neutron star radius is 10^6 cm. All the quantities in equation (13) are normalized to their typical values. One can see that the scattering efficiency may be as large as about a few tens. To conclude whether the condition of efficient intensity transfer given by equation (12) can indeed be satisfied let us estimate the level of background radiation resulting from the spontaneous scattering of the pulsar radio beam, $I^{(0)}_{\nu_b} \sim I^{(0)}_{\nu_a} n_e \eta r \mathrm{d}\sigma/\mathrm{d}\Omega_1$. Taking into account that $\mathrm{d}\sigma/\mathrm{d}\Omega_1 = r_e^2\sin^2\theta\sin^2\theta_1/(\gamma^6\eta_1^4\eta^2)$, $\eta \approx \theta^2/2$ and $\eta_1 \approx 1/\gamma^2$, we find that

$$\frac{I^{(0)}_{\nu_b}}{I^{(0)}_{\nu_a}} \sim 2n_e r_e^2 r \sim 10^{-10}P^{-1}\frac{\kappa}{10^2}\frac{B_\star}{10^{12}\,\mathrm{G}}\left(\frac{r}{10^8\,\mathrm{cm}}\right)^{-2}. \quad (14)$$

Hence, typically $I^{(0)}_{\nu_b}/I^{(0)}_{\nu_a} \sim 10^{-8} - 10^{-12}$, and to satisfy equation (12) $\Gamma = 18 - 28$ are necessary. Thus, the intensity transfer from the radio beam to the background can indeed be significant in pulsars, especially at low enough frequencies.

Let us consider the location of the scattering region in the magnetosphere of a pulsar. According to equation (13), Γ shows strong explicit dependence on r, $\Gamma \propto r^{-4}$. However, the implicit dependence on r appears still more significant. As $\theta \propto r$ (see below), $\nu_a = \nu_b/\theta^2\gamma^2 \propto r^{-2}$. This can be understood as follows. At a fixed frequency ν_b, the background radiation actually grows because of the scattering of different radio beam frequencies ν_a, which satisfy the condition $\nu_b = \nu_a\theta^2(r)\gamma^2$ at different altitudes r. Larger altitudes imply lower frequencies ν_a, in which case the incident intensity of the radio beam is much larger and stimulates stronger scattering. Taking into account the above considerations, one can obtain that $\Gamma \propto r^{2\alpha-4}$. Hence, the scattering efficiency increases with distance at $\alpha > 2$ and decreases at $\alpha < 2$. At high enough frequencies, $\gtrsim 1$ GHz, three of the pulsars with the precursor components, namely the Crab, the Vela and PSR B1822-09, have $\alpha = 2.8$, 2.7 and 2.3, respectively, whereas for PSR B1055-52 there is no spectral data in this region (perhaps, the spectrum is too steep for the pulsar to be detectable). At lower frequencies such a steep spectrum is preserved only in the Crab pulsar, while in the other pulsars under consideration α drops below 2. Note, however, that at lower frequencies the scattering is more efficient and may noticeably suppress the radio beam intensity, leading to the pulsar spectrum flattening, so that the original spectra may be much steeper than the observed ones. Thus, we assume that Γ increases with altitude and, correspondingly, the induced scattering is most efficient at the upper boundary of the scattering region, i.e. at distances of order of the cyclotron resonance radius, which is defined as $2\pi\nu_a\gamma\theta^2/2 = \omega_G$ and estimated as

$$\frac{r_c}{r_L} = 0.4\left(\frac{B_\star}{10^{12}\,\mathrm{G}}\frac{10^2}{\gamma}\frac{10^8\,\mathrm{Hz}}{\nu_a}\right)^{1/5}\left(\frac{1\,\mathrm{s}}{P}\right)^{3/5}, \quad (15)$$

where r_L is the light cylinder radius and it is taken that $\theta = r/2r_L$ (see below). One can see that the region of cyclotron resonance typically lies in the outer magnetosphere, at distances of order of the light cylinder radius.

The location of the scattered component in the pulse profile can be examined as follows. Since the scattering region lies well above the emission region, the wavevector $\boldsymbol{k}$ makes the angle $\sim r\sin\zeta/r_L$ with the instantaneous direction of the magnetic axis of the rotating magnetosphere (here ζ is the angle between the rotational and magnetic axes of a pulsar). As the polar angle of the point of scattering is $r\sin\zeta/r_L$, in the dipolar geometry the local magnetic field vector $\boldsymbol{b}$ is inclined at the angle $3r\sin\zeta/2r_L$ to the magnetic axis, so that the angle between $\boldsymbol{k}$ and $\boldsymbol{b}$ is $r\sin\zeta/2r_L$. In the corotating frame, the scattered radiation is directed approximately along $\boldsymbol{b}$. Then in the laboratory frame it is shifted because of rotational aberration by the angle $r\sin\zeta/r_L$ in the direction of rotation, i.e. toward the magnetic axis. Thus, the scattered component precedes the main pulse by $\Delta\lambda \sim r\sin\zeta/2r_L$ in longitude. One can see that the precursor separation from the main pulse is determined by the height of the scattering region, r, and does not exceed $\sim 30°$ for the scattering inside the light cylinder.

5 DISCUSSION

We have considered the induced Compton scattering by the particles of the ultrarelativistic electron-positron plasma in the presence of a superstrong magnetic field. In particular, we have examined the scattering of pulsar radio beam into background, which takes place in the open field line tube of a pulsar. It has been demonstrated that the photons are predominantly scattered approximately along the ambient magnetic field. This contrasts with the non-magnetic scattering, in which case the scattered photons concentrate in the backward direction. This difference is solely determined by a specific role of the superstrong magnetic field in the scattering process and does not depend on a detailed form of the particle distribution function.

Induced scattering in a superstrong magnetic field transfers the photons from lower to higher frequencies, $\nu_b \sim \nu_a\theta^2\gamma^2 \sim n \cdot 10\nu_a$, and if the process is efficient, the scattered component may become as strong as the original radio beam, $I_{\nu_b}(\nu_b) \sim I_{\nu_a}^{(0)}(\nu_a)$. As the beam has a decreasing spectrum, $I_{\nu_a}^{(0)}(\nu_a) \gg I_{\nu_a}^{(0)}(\nu_b)$, the intensity of the scattered component may dominate the original beam intensity at the same frequency ν_b.

For steep enough original spectra of pulsar radiation, $\alpha > 2$, the induced scattering in a superstrong magnetic field is most efficient at distances roughly comparable to the radius of cyclotron resonance. Because of rotational aberration, the scattered component appears in the pulse profile as a precursor to the main pulse. This effect provides the main pulse-precursor separations in longitude $\Delta\lambda \sim r\sin\zeta/2r_L$, which may run up to $\sim 30^\circ$. Since the length of the scattering region is larger than the height of the emission region, the intrinsic radius-to-frequency mapping of the radio emission is smeared. The effective height of the scattering region is an extremely weak function of the wave frequency, so that the main pulse-precursor separation is practically independent of frequency, just as is observed.

Since the induced scattering in the superstrong magnetic field holds only between the ordinary waves, the scattered component should have complete linear polarization. This is indeed the main distinctive feature of the observed precursors. Note that in general $[\boldsymbol{k}\times\boldsymbol{b}] \not\parallel [\boldsymbol{k}_1\times\boldsymbol{b}]$, i.e. in the main pulse and precursor the position angles of linear polarization should somewhat differ. Such a difference can be noticed, e.g., in PSR B1822-09 (Fowler et al. 1981). Besides that, if the main pulse is dominated by the extraordinary rather than ordinary polarization, the position angle of the precursor should additionally differ by 90°, as is the case in the Vela pulsar (Krishnamohan & Downs 1983).

As first noted by Fowler et al. (1981), the precursor components are met in pulsars with relatively large surface magnetic field. Firstly, large $B_\star$ are necessary for the regime of superstrong magnetic field to hold well above the emission region. Secondly, the scattering efficiency is proportional to $B_\star$. Short periods and large radio luminosities also favour significant scattering.

The pulse-to-pulse variations of the incident intensity and of the physical parameters in the scattering region may result in strong fluctuations of the precursor emission. The former variations imply the main pulse-precursor connection, which may have diversiform observational manifestations. For example, PSR J1326-6700 shows occasional main pulse nullings accompanied by the strong precursor emission (Wang et al. 2007). This can be interpreted as a consequence of extremely strong scattering, $(I_{\nu_b}^{(0)}/I_{\nu_a}^{(0)})\exp(\Gamma) \gg 1$, when the main pulse intensity is almost completely transferred to the precursor, $I_{\nu_a} \to 0$, $I_{\nu_b} \to I_{\nu_a}^{(0)}$. Note that this may happen only if the original intensity is mainly in the ordinary mode, which is subject to the scattering. In case of a moderately strong scattering, $(I_{\nu_b}^{(0)}/I_{\nu_a}^{(0)})\exp(\Gamma) \sim 1$, the main pulse intensity $I_{\nu_a}^{(0)}$ is almost unchanged, whereas the precursor grows exponentially with $I_{\nu_a}^{(0)}$. Therefore even weak fluctuations of the latter quantity may affect the scattered component dramatically. In some pulsars the precursors are indeed met only in strong pulses (Hankins & Cordes 1981; Gil et al. 1994; Weltevrede et al. 2006), and one can expect that the transient precursors are much more abundant in the pulsar population and are yet to be studied observationally.

The precursor component can fluctuate not only in intensity but also in pulse longitude. In the Vela pulsar, stronger precursors exhibit larger separations from the main pulse, which is thought to result from the fluctuations of the physical parameters in the scattering region. Larger separations imply larger scattering heights, $\Delta\lambda \propto r$, in which case the angle of incidence of the photons is also larger, $\theta \propto r$, and at a fixed frequency ν_b the precursor is formed by the photons coming from lower frequencies $\nu_a = \nu_b/\theta^2(r)\gamma^2$, which are more numerous and stimulate stronger scattering.

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